Scalar tetraquarks with open bottom

S.M. Gerasyuta * and V.I. Kochkin

Department of Physics, LTA, 194021, St. Petersburg, Russia

* E-mail: gerasyuta@SG6488.spb.edu

Abstract

The relativistic four-quark equations in the framework of the dispersion relation technique are derived. The four-quark amplitudes of the scalar tetraquarks with open bottom, including u, d, s and bottom quarks, are constructed. The poles of these amplitudes determine the masses and widths of scalar tetraquarks.

Keywords: scalar tetraquarks with open bottom, coupled-channel formalism.

PACS number: 11.55.Fv, 12.39.Ki, 12.39.Mk, 12.40.Yx.

I. Introduction.

The observation of the $X(3872)$ [1 – 4], the first of the XYZ particles to be seen, brought forward the hope that a multiquark state has been received. Belle Collaboration discovered the $X(3940)$ in the reaction $e^+e^- \to J/\psi + X$ [5]. The fact that the newly found states do not fit quark model calculations [6] has triggered the interest in the charmonium-like and the charmed states. Maiani et al. advocate a tetraquark explanation for the $X(3872)$ [7, 8]. Ebert et al. [9] calculated that the light scalar tetraquark lies above the open charm threshold and is broad, while Maiani et al. obtained [10] that this state lies a few MeV below this threshold. In our paper [11] the tetraquark $X(3700)$ with the spin-parity $J^{pc} = 0^{++}$ decays to $\eta\eta_c$. The calculated width of this state is equal to $\Gamma_{0^{++}} = 140\, MeV$.

In the present paper the relativistic four-quark equations are found in the framework of coupled-channel formalism. The dynamical mixing between the meson-meson states and the four-quark states is considered [12 – 14]. The masses and the widths of the scalar tetraquarks with open bottom are calculated.

II. Four-Quark Amplitudes for the Tetraquarks with Open Bottom.

We derive the relativistic four-quark equations in the framework of the dispersion relations technique.

The correct equations for the amplitude are obtained by taking into account all possible subamplitudes. It corresponds to the division of complete system into subsystems with the smaller number of particles. Then one should represent a four-particle amplitude as a sum of six subamplitudes:

$$A = A_{12} + A_{13} + A_{14} + A_{23} + A_{24} + A_{34}\,. \quad (1)$$

This defines the division of the diagrams into groups according to the certain pair interaction of particles. The total amplitude can be represented graphically as a sum of diagrams.

We need to consider only one group of diagrams and the amplitude corresponding, for example A_{12}.

The relativistic generalization of the Faddeev-Yakubovsky approach [15, 16] for the tetraquark is obtained. We shall construct the four-quark amplitude of $\bar{b}u\bar{u}u$ tetraquark with the spin-parity $J^{pc} = 0^{++}$ in which the quark amplitudes with quantum numbers of 0^{-+} and 1^{--} mesons are included. The set of diagrams associated with the amplitude A_{12} can further broken down into four groups corresponding to subamplitudes: $A_1(s, s_{12}, s_{34})$, $A_2(s, s_{23}, s_{14})$, $A_3(s, s_{23}, s_{123})$, $A_4(s, s_{14}, s_{124})$, if we consider the tetraquark with the spin-parity $J^{pc} = 0^{++}$ $(\bar{b}u\bar{u}u)$.

Here s_{ik} is the two-particle subenergy squared, s_{ijk} corresponds to the energy squared of particles i, j, k and s is the system total energy squared.

In order to represent the subamplitudes $A_1(s, s_{12}, s_{34})$, $A_2(s, s_{23}, s_{14})$, $A_3(s, s_{23}, s_{123})$ and $A_4(s, s_{14}, s_{124})$ in the form of dispersion relations it is necessary to define the amplitudes of quark-antiquark interaction $a_n(s_{ik})$. The pair quarks amplitudes $q\bar{q} \to q\bar{q}$ are calculated in the framework of the dispersion N/D method with the input four-fermion interaction [17 – 19] with quantum numbers of the gluon [20]. The regularization of the dispersion integral for the D-function is carried out with the cutoff parameter Λ. The four-quark interaction is considered as an input [20]:

$$g_V \left(\bar{q}\lambda I_f \gamma_\mu q\right)^2 + g_V^{(s)} \left(\bar{q}\lambda I_f \gamma_\mu q\right)\left(\bar{s}\lambda\gamma_\mu s\right) + g_V^{(ss)} \left(\bar{s}\lambda\gamma_\mu s\right)^2 . \tag{2}$$

Here I_f is the unity matrix in the flavor space (u, d). λ are the color Gell-Mann matrices. Dimensional constants of the four-fermion interaction g_V, $g_V^{(s)}$ and $g_V^{(ss)}$ are parameters of the model. At $g_V = g_V^{(s)} = g_V^{(ss)}$ the flavor $SU(3)_f$ symmetry occurs. The strange quark violates the flavor $SU(3)_f$ symmetry. In order to avoid an additional violation parameters, we introduce the scale shift of the dimensional parameters [20]:

$$g = \frac{m^2}{\pi^2} g_V = \frac{(m+m_s)^2}{4\pi^2} g_V^{(s)} = \frac{m_s^2}{\pi^2} g_V^{(ss)} . \tag{3}$$

$$\Lambda = \frac{4\Lambda(ik)}{(m_i + m_k)^2}. \tag{4}$$

Here m_i and m_k are the quark masses in the intermediate state of the quark loop. Dimensionless parameters g and Λ are supposed to be constants which are independent of the quark interaction type. The applicability of Eq. (2) is verified by the success of De Rujula-Georgi-Glashow quark model [21], where only the short-range part of Breit potential connected with the gluon exchange is responsible for the mass splitting in hadron multiplets.

We use the results of our relativistic quark model [20] and write down the pair quarks amplitude in the form:

$$a_n(s_{ik}) = \frac{G_n^2(s_{ik})}{1 - B_n(s_{ik})} , \tag{5}$$

$$B_n(s_{ik}) = \int\limits_{(m_i+m_k)^2}^{\frac{(m_i+m_k)^2\Lambda}{4}} \frac{ds'_{ik}}{\pi} \frac{\rho_n(s'_{ik}) G_n^2(s'_{ik})}{s'_{ik} - s_{ik}} . \tag{6}$$

Here $G_n(s_{ik})$ are the quark-antiquark vertex functions. The vertex functions are determined by the contribution of the crossing channels. The vertex functions satisfy

the Fierz relations. All of these vertex functions are generated from g_V, $g_V^{(s)}$ and $g_V^{(ss)}$. $B_n(s_{ik})$, $\rho_n(s_{ik})$ are the Chew-Mandelstam functions with cutoff Λ and the phase spaces, respectively.

Here $n = 1$ determines a $q\bar{q}$-pairs with $J^{pc} = 0^{-+}$ in the 1_c color state, $n = 2$ corresponds to a $q\bar{q}$-pairs with $J^{pc} = 1^{--}$ in the 1_c color state, and $n = 3$ defines the $q\bar{q}$-pairs corresponding to tetraquarks with quantum numbers $J^{pc} = 0^{++}$.

In the case in question, the interacting quarks do not produce a bound state; therefore, the integration in Eqs. (7) – (10) is carried out from the threshold $(m_i + m_k)^2$ to the cutoff $\Lambda(ik)$. The coupled integral equation systems (the tetraquark state with $n = 3$ and $J^{pc} = 0^{++}$ $\bar{b}u\bar{u}u$) can be described as:

$$\begin{aligned} A_1(s, s_{12}, s_{34}) &= \frac{\lambda_1 B_2(s_{12}) B_2(s_{34})}{[1 - B_2(s_{12})][1 - B_2(s_{34})]} + 2\hat{J}_2(s_{12}, s_{34}, 2, 2) A_3(s, s'_{23}, s'_{123}) \\ &+ 2\hat{J}_2(s_{12}, s_{34}, 2, 2) A_4(s, s'_{14}, s'_{124}) \,, \end{aligned} \tag{7}$$

$$\begin{aligned} A_2(s, s_{23}, s_{14}) &= \frac{\lambda_2 B_1(s_{23}) B_1(s_{14})}{[1 - B_1(s_{23})][1 - B_1(s_{14})]} + 2\hat{J}_2(s_{23}, s_{14}, 1, 1) A_3(s, s'_{34}, s'_{234}) \\ &+ 2\hat{J}_2(s_{23}, s_{14}, 1, 1) A_4(s, s'_{12}, s'_{123}) \,, \end{aligned} \tag{8}$$

$$\begin{aligned} A_3(s, s_{23}, s_{123}) &= \frac{\lambda_3 B_3(s_{23})}{[1 - B_3(s_{23})]} + 2\hat{J}_3(s_{23}, 3) A_1(s, s'_{12}, s'_{34}) + \hat{J}_3(s_{23}, 3) A_2(s, s'_{12}, s'_{34}) \\ &+ \hat{J}_1(s_{23}, 3) A_4(s, s'_{34}, s'_{234}) + \hat{J}_1(s_{23}, 3) A_3(s, s'_{12}, s'_{123}) \,, \end{aligned} \tag{9}$$

$$\begin{aligned} A_4(s, s_{14}, s_{124}) &= \frac{\lambda_4 B_3(s_{14})}{[1 - B_3(s_{14})]} + 2\hat{J}_3(s_{14}, 3) A_1(s, s'_{13}, s'_{24}) + 2\hat{J}_3(s_{14}, 3) A_2(s, s'_{13}, s'_{24}) \\ &+ 2\hat{J}_1(s_{14}, 3) A_3(s, s'_{14}, s'_{134}) + 2\hat{J}_1(s_{14}, 3) A_4(s, s'_{14}, s'_{134}) \,, \end{aligned} \tag{10}$$

where λ_i, $i = 1, 2, 3, 4$ are the current constants. They do not affect the mass spectrum of tetraquarks. We introduce the integral operators:

$$\hat{J}_1(s_{12}, l) = \frac{G_l(s_{12})}{[1 - B_l(s_{12})]} \int\limits_{(m_1+m_2)^2}^{\frac{(m_1+m_2)^2 \Lambda}{4}} \frac{ds'_{12}}{\pi} \frac{G_l(s'_{12}) \rho_l(s'_{12})}{s'_{12} - s_{12}} \int\limits_{-1}^{+1} \frac{dz_1}{2} \,, \tag{11}$$

$$\begin{aligned} \hat{J}_2(s_{12}, s_{34}, l, p) &= \frac{G_l(s_{12}) G_p(s_{34})}{[1 - B_l(s_{12})][1 - B_p(s_{34})]} \int\limits_{(m_1+m_2)^2}^{\frac{(m_1+m_2)^2 \Lambda}{4}} \frac{ds'_{12}}{\pi} \frac{G_l(s'_{12}) \rho_l(s'_{12})}{s'_{12} - s_{12}} \\ &\times \int\limits_{(m_3+m_4)^2}^{\frac{(m_3+m_4)^2 \Lambda}{4}} \frac{ds'_{34}}{\pi} \frac{G_p(s'_{34}) \rho_p(s'_{34})}{s'_{34} - s_{34}} \int\limits_{-1}^{+1} \frac{dz_3}{2} \int\limits_{-1}^{+1} \frac{dz_4}{2} \,, \end{aligned} \tag{12}$$

$$\hat{J}_3(s_{12},l) = \frac{G_l(s_{12},\tilde{\Lambda})}{[1-B_l(s_{12},\tilde{\Lambda})]} \frac{1}{4\pi} \int\limits_{(m_1+m_2)^2}^{\frac{(m_1+m_2)^2\tilde{\Lambda}}{4}} \frac{ds'_{12}}{\pi} \frac{G_l(s'_{12},\tilde{\Lambda})\rho_l(s'_{12})}{s'_{12}-s_{12}}$$

$$\times \int\limits_{-1}^{+1} \frac{dz_1}{2} \int\limits_{-1}^{+1} dz \int\limits_{z_2^-}^{z_2^+} dz_2 \frac{1}{\sqrt{1-z^2-z_1^2-z_2^2+2zz_1z_2}}\,, \qquad (13)$$

here l, p are equal to $1-3$.

In Eqs. (11) and (13) z_1 is the cosine of the angle between the relative momentum of particles 1 and 2 in the intermediate state and the momentum of the particle 3 in the final state, taken in the c.m. of particles 1 and 2. In Eq. (13) z is the cosine of the angle between the momenta of particles 3 and 4 in the final state, taken in the c.m. of particles 1 and 2. z_2 is the cosine of the angle between the relative momentum of particles 1 and 2 in the intermediate state and the momentum of the particle 4 in the final state, is taken in the c.m. of particles 1 and 2. In Eq. (12): z_3 is the cosine of the angle between the relative momentum of particles 1 and 2 in the intermediate state and the relative momentum of particles 3 and 4 in the intermediate state, taken in the c.m. of particles 1 and 2. z_4 is the cosine of the angle between the relative momentum of particles 3 and 4 in the intermediate state and that of the momentum of the particle 1 in the intermediate state, taken in the c.m. of particles 3 and 4.

We can pass from the integration over the cosines of the angles to the integration over the subenergies [22].

Let us extract two-particle singularities in the amplitudes $A_1(s,s_{12},s_{34})$, $A_2(s,s_{23},s_{14})$, $A_3(s,s_{23},s_{123})$ and $A_4(s,s_{14},s_{124})$:

$$A_1(s,s_{ik},s_{lm}) = \frac{\alpha_1(s,s_{ik},s_{lm})B_2(s_{ik})B_2(s_{lm})}{[1-B_2(s_{ik})][1-B_2(s_{lm})]}\,, \qquad (14)$$

$$A_2(s,s_{ik},s_{lm}) = \frac{\alpha_2(s,s_{ik},s_{lm})B_1(s_{ik})B_1(s_{lm})}{[1-B_1(s_{ik})][1-B_1(s_{lm})]}\,, \qquad (15)$$

$$A_j(s,s_{ik},s_{ikl}) = \frac{\alpha_j(s,s_{ik},s_{ikl})B_3(s_{ik})}{1-B_3(s_{ik})}\,, \qquad j=3-4\,. \qquad (16)$$

We do not extract three-particles singularities, because they are weaker than two-particle singularities.

We used the classification of singularities, which was proposed in paper [23]. The construction of the approximate solution of Eqs. (7) – (10) is based on the extraction of the leading singularities of the amplitudes. The main singularities in $s_{ik} \approx (m_i+m_k)^2$ are from pair rescattering of the particles i and k. First of all there are threshold square-root singularities. Also possible are pole singularities which correspond to the bound states. The amplitudes apart from two-particle singularities have triangular singularities and the singularities defining the interactions of four particles. Such classification allows us to search the corresponding solution of Eqs. (7) – (10) by taking into account some definite number of leading singularities and neglecting all the weaker ones. We consider the approximation which defines two-particle, triangle and four-particle singularities. The functions $\alpha_1(s,s_{12},s_{34})$, $\alpha_2(s,s_{23},s_{14})$, $\alpha_3(s,s_{23},s_{123})$ and $\alpha_4(s,s_{14},s_{124})$ are the smooth functions of s_{ik}, s_{ikl}, s as compared with the singular part of the amplitude, hence they can be expanded in a series in the singulary point and only the first term of this series should be employed further. Using this classification, one defines

the reduced amplitudes α_1, α_2, α_3, α_4 as well as the B-functions in the middle point of physical region of Dalitz-plot at the point s_0:

$$s_0^{ik} = 0.25(m_i + m_k)^2 s_0 \,,$$

$$s_{123} = 0.25 s_0 \sum_{\substack{i,k=1\\ i\neq k}}^{3} (m_i + m_k)^2 - \sum_{i=1}^{3} m_i^2 \,, \quad s_0 = \frac{s + 2\sum\limits_{i=1}^{4} m_i^2}{0.25 \sum\limits_{\substack{i,k=1\\ i\neq k}}^{4} (m_i + m_k)^2} \,. \tag{17}$$

Such a choice of point s_0 allows us to replace integral Eqs. (7) – (10) by the algebraic equations (18) – (21) respectively:

$$\alpha_1 = \lambda_1 + 2\alpha_3 JB_1(2,2,3) + 2\alpha_4 JB_2(2,2,3) \,, \tag{18}$$

$$\alpha_2 = \lambda_2 + 2\alpha_3 JB_3(1,1,3) + 2\alpha_4 JB_4(1,1,3) \,, \tag{19}$$

$$\alpha_3 = \lambda_3 + 2\alpha_1 JC_1(3,2,2) + 2\alpha_2 JC_2(3,1,1) + \alpha_4 JA_1(3) + \alpha_3 JA_2(3) \,, \tag{20}$$

$$\alpha_4 = \lambda_4 + 2\alpha_1 JC_3(3,2,2) + 2\alpha_2 JC_4(3,1,1) + \alpha_3 JA_3(3) + \alpha_4 JA_4(3) \,. \tag{21}$$

We use the functions $JA_i(l)$, $JB_i(l,p,r)$, $JC_i(l,p,r)$ $(l,p,r = 1-3)$, which are determined by the various s_0^{ik} (Eq. 17). These functions are similar to the functions:

$$JA_4(l) = \frac{G_l^2(s_0^{12}) B_l^2(s_0^{23})}{B_l(s_0^{12})} \int\limits_{(m_1+m_2)^2}^{\frac{(m_1+m_2)^2\Lambda}{4}} \frac{ds'_{12}}{\pi} \frac{\rho_l(s'_{12})}{s'_{12} - s_{12}} \int\limits_{-1}^{+1} \frac{dz_1}{2} \frac{1}{1 - B_l(s'_{23})} \,, \tag{22}$$

$$\begin{aligned} JB_1(l,p,r) &= \frac{G_l^2(s_0^{12}) G_p^2(s_0^{34}) B_r(s_0^{23})}{B_l(s_0^{12}) B_p(s_0^{34})} \int\limits_{(m_1+m_2)^2}^{\frac{(m_1+m_2)^2\Lambda}{4}} \frac{ds'_{12}}{\pi} \frac{\rho_l(s'_{12})}{s'_{12} - s_{12}} \\ &\times \int\limits_{(m_3+m_4)^2}^{\frac{(m_3+m_4)^2\Lambda}{4}} \frac{ds'_{34}}{\pi} \frac{\rho_p(s'_{34})}{s'_{34} - s_{34}} \int\limits_{-1}^{+1} \frac{dz_3}{2} \int\limits_{-1}^{+1} \frac{dz_4}{2} \frac{1}{1 - B_r(s'_{23})} \,, \end{aligned} \tag{23}$$

$$\begin{aligned} JC_3(l,p,r) &= \frac{G_l^2(s_0^{12}, \tilde{\Lambda}) B_p(s_0^{23}) B_r(s_0^{14})}{1 - B_l(s_0^{12}, \tilde{\Lambda})} \frac{1 - B_l(s_0^{12})}{B_l(s_0^{12})} \frac{1}{4\pi} \int\limits_{(m_1+m_2)^2}^{\frac{(m_1+m_2)^2\tilde{\Lambda}}{4}} \frac{ds'_{12}}{\pi} \frac{\rho_l(s'_{12})}{s'_{12} - s_{12}} \\ &\times \int\limits_{-1}^{+1} \frac{dz_1}{2} \int\limits_{-1}^{+1} dz \int\limits_{z_2^-}^{z_2^+} dz_2 \frac{1}{\sqrt{1 - z^2 - z_1^2 - z_2^2 + 2zz_1z_2}} \\ &\times \frac{1}{[1 - B_p(s'_{23})][1 - B_r(s'_{14})]} \,, \end{aligned} \tag{24}$$

$$\tilde{\Lambda}(ik) = \begin{cases} \Lambda(ik), & \text{if} \quad \Lambda(ik) \leq (\sqrt{s_{123}} + m_3)^2 \\ (\sqrt{s_{123}} + m_3)^2, & \text{if} \quad \Lambda(ik) > (\sqrt{s_{123}} + m_3)^2 \end{cases} \tag{25}$$

The other choices of point s_0 do not change essentially the contributions of α_1, α_2, α_3 and α_4, therefore we omit the indices s_0^{ik}. Since the vertex functions depend only slightly on energy it is possible to treat them as constants in our approximation.

The solutions of the system of equations are considered as:

$$\alpha_i(s) = F_i(s, \lambda_i)/D(s)\,, \tag{26}$$

where zeros of $D(s)$ determinants define the masses of bound states of tetraquarks. $F_i(s, \lambda_i)$ determine the contributions of subamplitudes to the tetraquark amplitude.

III. Calculation results.

Our calculations do not include the new parameters. We use the cutoff $\Lambda = 7.63$ and the gluon coupling constant $g = 1.53$, which are determined by fixing the tetraquark masses for the states with the hidden bottom [24, 25]. The widths of the tetraquarks with the open bottom are fitted by the fixing width $\Gamma_{2^{++}} = (39 \pm 26)\,MeV$ [26] for the S-wave tetraquark with the hidden charm $X(3940)$. The quark masses of model $m_{u,d} = 385\,MeV$, $m_s = 510\,MeV$ and $m_b = 4787\,MeV$ coincide with our paper ones [25]. The masses and the widths of meson-meson states with the spin-parity $J^{pc} = 0^{++}$ are given in Table I. In our paper we predicted the tetraquark $(\bar{b}u\bar{u}u)$ with the mass $M = 5914\,MeV$ and the width $\Gamma_{0^{++}} = 104\,MeV$. We calculated the masses of $X_b(6020)$ $M = 6017\,MeV$ and the width $\Gamma_{0^{++}} = 69\,MeV$ (channels $B_s^0\eta$ and B^+K^-). The tetraquark $(\bar{b}su\bar{s})$ have the mass $M = 6122\,MeV$ and the width $\Gamma_{0^{++}} = 48\,MeV$. The tetraquarks with the open bottom and the spin-parity $J^{pc} = 1^{++}$, 2^{++} have only the weak decays.

The functions $F_i(s, \lambda_i)$ (Eq. (26)) allow us to obtain the overlap factors f for the tetraquarks. We calculated the overlap factors f and the phase spaces ρ for the reactions $X \to M_1 M_2$ (Table I). We considered the formula $\Gamma \sim f^2 \times \rho$ [27], there ρ is the phase space.

In the open bottom sector the scalar tetraquarks have relatively small width $\sim 50 - 100\,MeV$, so in principle, these exotic states could be observed.

Acknowledgments.

The work was carried with the support of the Russian Ministry of Education (grant 2.1.1.68.26).

Table I. Masses, widths, overlap factors f and phase spaces ρ of scalar tetraquarks with open bottom.

Tetraquark	(channels)	f	ρ	Mass (MeV)	Widths (MeV)
$X_b(5910)$	$B^+\eta$	0.54	0.103	5914	104
$X_b(6020)$	$B_s^0\eta$ B^+K^-	0.32 0.23	0.108 0.171	6017	69
$X_b(6120)$	$B_s^0K^+$	0.283	0.174	6122	48

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